arXiv:0712.2519
YITP-07-90
KUNS-2116

Metastable supersymmetry breaking vacua from conformal dynamics

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Abstract

We study the scenario that conformal dynamics leads to metastable supersymmetry breaking vacua. At a high energy scale, the superpotential is not R-symmetric, and has a supersymmetric minimum. However, conformal dynamics suppresses several operators along renormalization group flow toward the infrared fixed point. Then we can find an approximately R-symmetric superpotential, which has a metastable supersymmetry breaking vacuum, and the supersymmetric vacuum moves far away from the metastable supersymmetry breaking vacuum. We show a 4D simple model. Furthermore, we can construct 5D models with the same behavior, because of the AdS/CFT dual.

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1 Introduction

Conformal dynamics provides several interesting aspects in supersymmetric models as well as non-supersymmetric models, because conformal dynamics exponentially suppresses or enhances certain operators. For example, contact terms like $\int d^4\theta |X|^2 |Q|^2$ are suppressed exponentially by conformal dynamics in the model that the chiral superfield X belongs to the hidden conformal sector and the chiral superfield Q belongs not to the conformal sector, but to the visible sector. Such conformal suppression mechanism, i.e. conformal sequestering, is quite important to model building for supersymmetry (SUSY) breaking [1, 2, 3, 4, 5, 6]. When X contributes to SUSY breaking sizably, the above contact terms, in general, induce flavor-dependent soft SUSY breaking terms, soft sfermion masses and the so-called A-terms, and they lead to flavor changing neutral current processes, which are strongly constrained by current experiments. However, conformal sequestering can suppress the above contact terms and flavor-dependent contributions to soft SUSY breaking terms. Then, flavor-blind contributions such as anomaly mediation [7] would become dominant.

Another interesting aspect is that conformal dynamics can generate hierarchical structure of Yukawa couplings for quarks and leptons [8, 9]. Suppose that our quark and lepton superfields Q couple with the conformal sector like $\int d^2\theta h Q X_1 X_2$, and the fields X_1 and X_2 have negative anomalous dimensions by conformal dynamics. When this coupling h is driven toward an infra-red (IR) fixed point, the fields Q would have large and positive anomalous dimensions. Because of such large and positive anomalous dimensions, Yukawa couplings among quarks/leptons Q and the electroweak Higgs fields become exponentially suppressed toward the IR direction. Then, even if all of Yukawa couplings are of $O(1)$ at a high energy scale, hierarchies among Yukawa couplings could be generated by conformal dynamics. At the same time, sfermion masses are exponentially suppressed toward the IR fixed point [8, 9, 10].[1]

Here we study a new application of conformal dynamics for supersymmetric models, that is, realization of metastable SUSY breaking vacua by conformal dynamics. Its idea is as follows. The Nelson-Seiberg argument [12] implies that generic superpotential has a SUSY minimum, but R-symmetric

[1] A similar dynamics would be useful to control a large radiative correction on Higgs soft masses [11].

superpotential has no SUSY minimum, that is, SUSY is broken in such a model. Thus, if we add explicit R-symmetry breaking terms in R-symmetric superpotential, a SUSY minimum would appear. However, when such R-symmetry breaking terms are tiny, the previous SUSY breaking minimum would survive and a newly appeared SUSY preserving minimum would be far away from the SUSY breaking point in the field space. That is the metastable SUSY breaking vacuum [13, 14, 15]. We try to realize such a metastable SUSY breaking vacuum by conformal dynamics. We start with a superpotential without R-symmetry. However, we assume the conformal dynamics. Because of that, certain couplings are exponentially suppressed. Then, we could realize an R-symmetric superpotential or an approximately R-symmetric superpotential with tiny R-symmetry breaking terms. It would lead to a stable or metastable SUSY breaking vacuum. We study this scenario by using a simple model. Also, we study 5D models, which have the same behavior.

This paper is organized as follows. In the next section, we give a 4D simple model to realize our conformal scenario. In section 3, we study 5D models, which have the same behavior. Section 4 is devoted to conclusion and discussion.

2 4D conformal model

Our model is the $SU(N)$ gauge theory with N_f flavors of chiral matter fields ϕ_i and $\tilde{\phi}_i$, which are fundamental and anti-fundamental representations of $SU(N)$. The flavor number satisfies $N_f \geq \frac{3}{2}N$, and that corresponds to the conformal window [16, 17], that is, this theory has an IR fixed point [18]. The NSVZ beta-function of physical gauge coupling $\alpha = g^2/8\pi^2$ is

$$\beta_\alpha^{\rm NSVZ} = -\frac{\alpha^2}{1-N\alpha}(3N - N_f + N_f\gamma_\phi), \tag{1}$$

where γ_ϕ is the anomalous dimension of ϕ_i and $\tilde{\phi}_i$ [19, 20]. Since the IR fixed point corresponds to $\beta_\alpha^{\rm NSVZ} = 0$, around that point the matter fields ϕ_i and $\tilde{\phi}_i$ have anomalous dimensions $\gamma_\phi = -(3N - N_f)/N_f$, which are negative.

In addition to the fields ϕ_i and $\tilde{\phi}_i$, we introduce singlet fields Φ_{ij} for $i, j = 1, \cdots, N_f$. The gauge invariance allows the following superpotential at

the renormalizable level,

$$W = h\phi_i \Phi_{ij} \tilde{\phi}_j + f \mathrm{Tr}_{ij} \Phi_{ij} + \frac{m}{2} \mathrm{Tr}_{ik} \Phi_{ij} \Phi_{jk} + \frac{\lambda}{3} \mathrm{Tr}_{i\ell} \Phi_{ij} \Phi_{jk} \Phi_{k\ell}. \tag{2}$$

Here we have preserved the $SU(N_f)$ flavor symmetry. Even if the $SU(N_f)$ flavor symmetry is broken, e.g by replacing $f \mathrm{Tr}_{ij} \Phi_{ij}$ by $f_{ij} \Phi_{ij}$, the following discussions would be valid. For simplicity, we assume that all of couplings, h, f, m, λ, are real, although the following discussions are available for the model with complex parameters, h, f, m and λ. We can add the mass terms of ϕ_i and $\tilde{\phi}_j$ to the above superpotential. We will comment on such terms later, but at the first stage we study the superpotential without the mass terms of ϕ_i and $\tilde{\phi}_j$.

If $m = \lambda = 0$, the above superpotential corresponds to the superpotential of the Intriligator-Seiberg-Shih (ISS) model [21]. We consider that our theory is an effective theory with the cutoff Λ. We assume that dimensionless parameters h and λ are of $O(1)$ and dimensionful parameters f and m satisfy $f \approx m^2$ and $m \ll \Lambda$. We denote physical couplings as $\hat{h} = (Z_\phi Z_{\tilde{\phi}} Z_\Phi)^{-1/2} h$, $\hat{f}_{ij} = (Z_\Phi)^{-1/2} f_{ij}$, $\hat{m} = (Z_\Phi)^{-1} m$ and $\hat{\lambda} = (Z_\Phi)^{-3/2} \lambda$, where $Z_\phi, Z_{\tilde{\phi}}, Z_\Phi$ are wavefunction renormalization constants for $\phi, \tilde{\phi}, \Phi$, respectively.

The F-flat conditions are obtained as

$$\begin{aligned}
\partial_{\Phi_{ij}} W &= h\phi_i \tilde{\phi}_j + f\delta_{ij} + m\Phi_{ij} + \lambda \Phi_{jk} \Phi_{ki} = 0, && (3) \\
\partial_{\phi_i} W &= h\Phi_{ij} \tilde{\phi}_j = 0, && (4) \\
\partial_{\tilde{\phi}_j} W &= h\phi_i \Phi_{ij} = 0. && (5)
\end{aligned}$$

These equations have a supersymmetric solution for generic values of parameters, h, m, λ. To see such a supersymmetric solution, following [21] we decompose $\phi, \tilde{\phi}$ and Φ as

$$\Phi = \begin{pmatrix} Y & Z^T \\ \tilde{Z} & X \end{pmatrix}, \qquad \phi = \begin{pmatrix} \chi \\ \rho \end{pmatrix}, \qquad \tilde{\phi}^T = \begin{pmatrix} \tilde{\chi} \\ \tilde{\rho} \end{pmatrix}, \tag{6}$$

where Y, χ and $\tilde{\chi}$ are $N \times N$ matrices, X is an $(N_F - N) \times (N_F - N)$ matrix, Z, $\tilde{Z}$, ρ and $\tilde{\rho}$ are $(N_F - N) \times N$ matrices. Let us consider the slice with $Z = \tilde{Z} = \rho = 0$ in the field space, where the first derivatives of W reduce to

$$W_{\Phi_{ij}} = \begin{pmatrix} f\delta_{ij} + h\chi_i \tilde{\chi}_j + mY_{ji} + \lambda Y_{jk} Y_{ki} & 0 \\ 0 & f\delta_{ij} + mX_{ji} + \lambda X_{jk} X_{ki} \end{pmatrix}, \tag{7}$$

$$W_{\phi_i}^T = \begin{pmatrix} hY_{ij}\tilde{\chi}_j \\ 0 \end{pmatrix}, \qquad W_{\tilde{\phi}_j} = \begin{pmatrix} h\chi_i Y_{ij} \\ 0 \end{pmatrix}. \tag{8}$$

Here, we have used the same indices for Φ_{ij}, ϕ_i, $\tilde{\phi}_j$ and their submatrices. Thus, the fields X_{ij} and the others are decoupled in the F-flat conditions, $W_{\Phi_{ij}} = W_{\phi_i} = W_{\tilde{\phi}_j} = 0$. The F-flat condition $W_{\Phi_{ij}} = 0$ for X_{ij} has a solution as $X_{ij} = x_s \delta_{ij}$ with

$$x_s = \frac{-m \pm \sqrt{m^2 - 4f\lambda}}{2\lambda}. \tag{9}$$

The F-flat conditions $W_{\Phi_{ij}} = W_{\phi_i} = W_{\tilde{\phi}_j} = 0$ for Y_{ij}, χ_i and $\tilde{\chi}_j$ have the following solution,

$$f\delta_{ij} + h\chi_i\tilde{\chi}_j = 0, \qquad Y_{ij} = 0. \tag{10}$$

In addition, the D-flat conditions correspond to $|\chi_i| = |\tilde{\chi}_i|$.

There is another solution, $\chi_i = \tilde{\chi}_j = 0$ and $Y_{ij} = x_s\delta_{ij}$. However, only the above solution (10) survives at the IR region, as $\hat{m}$ and $\hat{\lambda}$ become to vanish as we will see later. Thus, we concentrate to the solution (10). At any rate, the superpotential (2) does not have R-symmetry, and there is a supersymmetric minimum.

The above aspect is the behavior of this model around the energy scale Λ. Now let us study the behavior around the IR region. We assume that the gauge coupling is around the IR fixed point, i.e. $\beta_\alpha \approx 0$, and that ϕ_i and $\tilde{\phi}_i$ have negative anomalous dimensions γ_ϕ. In addition, we assume that the physical Yukawa coupling $\hat{h}$ is driven toward IR fixed points. The beta-function of $\hat{h}$ is obtained as

$$\beta_{\hat{h}} = \hat{h}(\gamma_\phi + \gamma_{\tilde{\phi}} + \gamma_\Phi). \tag{11}$$

The condition of the fixed point leads to $2\gamma_\phi + \gamma_\Phi = 0$. Since $\gamma_\phi < 0$, we obtain a positive anomalous dimension for Φ_{ij}. Then, physical couplings, $\hat{f}$, $\hat{m}$ and $\hat{\lambda}$, are suppressed exponentially toward the IR direction as

$$\begin{aligned} \hat{f}(\mu) &= \left(\frac{\mu}{\Lambda}\right)^{\gamma_\Phi} \hat{f}(\Lambda), \qquad \hat{m}(\mu) = \left(\frac{\mu}{\Lambda}\right)^{2\gamma_\Phi} \hat{m}(\Lambda), \\ \hat{\lambda}(\mu) &= \left(\frac{\mu}{\Lambda}\right)^{3\gamma_\Phi} \hat{\lambda}(\Lambda). \end{aligned} \tag{12}$$

Thus, the mass parameter $\hat{m}$ and 3-point coupling $\hat{\lambda}$ are suppressed faster than $\hat{f}$. If we neglect $\hat{m}$ and $\hat{\lambda}$ but not $\hat{f}$, the above superpotential becomes

the superpotential of the ISS model, and there is a SUSY breaking minimum around $\Phi_{ij} = 0$ because of the rank condition.

Let us see more explicitly. We concentrate ourselves to the potential of the fields X_{ij}, because X_{ij} contribute to SUSY breaking in the ISS model. Furthermore, we consider their overall direction, i.e. $X_{ij} = x\delta_{ij}$, and we use the canonically normalized basis, $\hat{x}$. Then, the above superpotential (2) leads to the following scalar potential,

$$V_{\rm SUSY} = (N_f - N)|\hat{f} + \hat{m}\hat{x} + \hat{\lambda}\hat{x}^2|^2. \tag{13}$$

In addition, around $\hat{x} = 0$, SUSY is broken and that generates one-loop effective potential of $\hat{x}$. Around $\hat{x} = 0$, the mass term $m_x^2|\hat{x}|^2$ in the one-loop effective potential would be important. Hence, we analyze the potential, $V = V_{\rm SUSY} + m_x^2|\hat{x}|^2$, and we use m_x^2, which has been calculated in [21], i.e.

$$m_x^2 = \frac{\hat{h}^3\hat{f}}{8\pi^2}N(N_f - N)(\log 4 - 1). \tag{14}$$

Note that m_x^2 is suppressed toward the IR region like $\hat{f}$. We consider only the real part of $\hat{x}$. The stationary condition $\partial_{\hat{x}}V =$ is written as

$$(\hat{f} + \hat{m}\hat{x} + \hat{\lambda}\hat{x}^2)(\hat{m} + 2\hat{\lambda}\hat{x}) + m_x^2\hat{x} = 0. \tag{15}$$

At a high energy scale corresponding to $Z_\Phi = O(1)$, we have $|\hat{f}|, |\hat{m}|^2 \gg m_x^2$, because m_x^2 is smaller than $\hat{f}$ by a loop factor. The potential and the stationary condition are controlled by $|\hat{f}|, |\hat{m}|^2$, $\hat{\lambda}$, but not m_x. Thus, there is no (SUSY breaking) minimum around $x = 0$, but we have a supersymmetric minimum

$$\hat{x}_s = \frac{-\hat{m} \pm \sqrt{\hat{m}^2 - 4\hat{f}\hat{\lambda}}}{2\hat{\lambda}}. \tag{16}$$

However, toward the IR direction, $\hat{m}^2$ becomes suppressed faster than m_x^2. Then, the couplings $\hat{f}$ and m_x^2 are important in the potential. Around $\hat{x} = 0$, the stationary condition (15) becomes

$$\hat{f}\hat{m} + m_x^2\hat{x} + \cdots = 0, \tag{17}$$

that is, the stationary condition is satisfied with

$$\hat{x}_{sb} \approx -\frac{\hat{f}\hat{m}}{m_x^2}. \tag{18}$$

At this point, SUSY is broken, and this point becomes close to $\hat{x}_{sb} = 0$ toward the IR. Around $\hat{x} = 0$, the size of mass is estimated by m_x, because the other terms are suppressed. Hence, the SUSY breaking metastable vacuum corresponding to $\hat{x} \sim 0$ appears at the IR energy scale, where $\hat{m}^2 \ll m_x^2$. Moreover, the previous SUSY vacuum (16) moves to a point far away from the origin $\hat{x} = 0$, because it behaves like

$$\hat{x}_s = \frac{-\hat{m} \pm \sqrt{\hat{m}^2 - 4\hat{f}\hat{\lambda}}}{2\hat{\lambda}} \sim \left(\frac{\Lambda}{\mu}\right)^{\gamma_\Phi}. \tag{19}$$

Both breaking scales of the $SU(N)$ gauge symmetry and supersymmetry at the metastable SUSY breaking point $\hat{x} = 0$ are determined by $O(\hat{f}(\mu))$. Thus, such an energy scale is estimated as $\mu_{IR}^2 \sim \hat{f}(\mu_{IR})$, i.e.

$$\mu_{IR} \sim \left(\frac{\hat{f}(\Lambda)}{\Lambda^{\gamma_\Phi}}\right)^{1/(2-\gamma_\Phi)}, \tag{20}$$

and at this energy scale conformal renormalization group flow is terminated.

So far, we have assumed that the mass term of ϕ_i and $\tilde{\phi}_i$, $m_\phi \phi_i \tilde{\phi}_i$ vanishes. Here, we comment on the case with such terms. The physical mass $\hat{m}_\phi$ becomes enhanced as

$$\hat{m}_\phi(\mu) = \left(\frac{\mu}{\Lambda}\right)^{2\gamma_\phi} \hat{m}_\phi(\Lambda), \tag{21}$$

because of the negative anomalous dimension γ_ϕ. At $\mu \sim \hat{m}_\phi(\mu)$, the matter fields ϕ_i $\tilde{\phi}_i$ decouple and this theory removes away from the conformal window. Thus, if $\hat{m}_\phi(\mu) > \mu_{IR}$, the conformal renormalization group flow is terminated at $\mu_D \sim \hat{m}_\phi(\mu_D) = (\mu_D/\Lambda)^{2\gamma_\phi}\hat{m}_\phi(\Lambda)$.

We have studied the scenario that conformal dynamics leads to metastable SUSY breaking vacua. As an illustrating example of our idea, we have used the simple model. Our scenario could be realized by other models.

3 5D model

There would be an AdS dual to our conformal scenario. Indeed, we can construct simply various models within the framework of 5D orbifold theory.

Renormalization group flows in the 4D theory correspond to exponential profiles of zero modes like $e^{-c_i R y}$, where R is the radius of the fifth dimension,[2] y is the coordinate for the extra dimension, i.e. $y = [0, \pi]$ and c_i is a constant. The parameter c_i corresponds to anomalous dimension in the 4D theory, and each field would have a different constant c_i. In 4D theory, values of anomalous dimensions are constrained by concrete 4D dynamics. However, constants c_i do not have such strong constraints, although they would correspond to some charges. Hence, 5D models would have a rich structure and one could make model building rather simply. Here we show a simple 5D model. We consider the 5D theory, whose 5-th dimension is compactified on S^1/Z_2. Two fixed points on S^1/Z_2 correspond to $y = 0$ and $y = \pi$. We introduce three bulk fields X, ϕ_1, ϕ_2. They correspond to chiral multiplets of bulk hyper-multiplets and zero modes of their partners in hyper-multiplets X^c, ϕ_1^c, ϕ_2^c are projected out by the Z_2 orbifold projection. We assume that zero mode profiles of X, ϕ_1 and ϕ_2 behave along the y direction as $e^{-c_X R y}$, $e^{-c_1 R y}$ and $e^{-c_2 R y}$, respectively. We integrate y and obtain their kinetic term coefficients Y_i of 4D effective theory, that is, the field corresponding to the zero mode profile $e^{-c_i R y}$ has the following kinetic term coefficient [22, 23]

$$Y_i = \frac{1}{c_i}\left(1 - e^{-2c_i \pi R}\right). \tag{22}$$

In the limit $c_i \to 0$, Y_i becomes $2\pi R$. Their superpotential is not allowed in the bulk, but is allowed on the boundary.

Suppose that the following superpotential is allowed only on the $y = \pi$ boundary,

$$\int dy \delta(y - \pi) W^{(\pi)}, \tag{23}$$

$$\begin{aligned} W^{(\pi)} \;\; = \;\; & f e^{-c_X R y} X + m e^{-2c_X R y} X^2 + h e^{-3c_X R y} X^3 \\ & + m_{12} e^{-(c_1+c_2) R y} \phi_1 \phi_2 + m_2 e^{-2c_2 R y} \phi_2^2 \\ & + \sum_{i,j} h_{ij} e^{-(c_X + c_i + c_j) R y} X \phi_i \phi_j . \end{aligned} \tag{24}$$

Here we have assumed extra Z_2 symmetry, under which X has the even Z_2 charge and ϕ_1 and ϕ_2 have the odd Z_2 charge. That allows the mass term $m_{11}\phi^2$, but we have assumed it vanishes by the same reason as why we did

[2] We assume that the radion is stabilized.

not add the mass term $m_{ij}\phi_i\tilde{\phi}_j$ in the superpotential (2). We assume that $f \approx m^2 \approx m_{12}^2 \approx m_2^2$ and $h, h_{ij} = O(1)$. We take

$$c_1 = 0, \qquad c_2 = c_X, \tag{25}$$

and $c_X > 0$ with $c_X\pi R = O(1)$ and $e^{-c_X\pi R} \ll 1$. The 4D superpotential $\hat{W}$ becomes

$$\hat{W} = e^{-c_X\pi R}(fX + m_{12}\phi_1\phi_2 + h_{11}X\phi_1^2) + e^{-2c_X\pi R}\Delta W. \tag{26}$$

When we neglect ΔW, the superpotential W corresponds to the O'Raifeartaigh model [24], that is, SUSY is broken. Such a minimum is metastable and there is a SUSY minimum, when we take into account ΔW [15]. The O'Raifeartaigh model with the following superpotential,

$$W_0 = \hat{f}X + \hat{m}_{12}\phi_1\phi_2 + \hat{h}_{11}X\phi_1^2, \tag{27}$$

leads to the SUSY breaking minimum of scalar potential $V = |\hat{f}|^2$ at $\phi_1 = \phi_2 = 0$ and arbitrary X, that is, it has the pseudo-flat direction. One-loop effects lift up this pseudo-flat direction, and the field X has the mass m_X,

$$m_X^2 = O\left(\frac{1}{4\pi^2}\frac{\hat{f}^2\hat{h}_{11}^4}{\hat{m}_{12}^2}\right), \tag{28}$$

around $X = 0$. In the case with $h_{11} = O(1)$ in the superpotential (26), we would have a rather small mass m_X by the suppression factor $e^{-c_X\pi R}$. To have larger mass m_X, we can assume the following superpotential $W^{(0)}$ at $y = 0$ as

$$\int dy\delta(y)W^{(0)} = \int dy\delta(y)h_{11}^{(0)}X\phi_1^2. \tag{29}$$

In this case, the 4D superpotential becomes

$$\hat{W} = (h_{11}^{(0)} - h_{11}e^{-C_X\pi R})X\phi_1^2 + e^{-c_X\pi R}(fX + m_{12}\phi_1\phi_2) + e^{-2c_X\pi R}\Delta W. \tag{30}$$

This leads to the metastable SUSY breaking minimum around $X = 0$ and the field X can have a larger mass around $X = 0$ than the previous model, because the coupling $h_{11}^{(0)}$ has no suppression factor like $e^{-c_X\pi R}$. The SUSY breaking source F^X is quasi-localized around $y = 0$.

We can construct more various models for approximately R-symmetric superpotential with metastable SUSY breaking vacua in 5D theory. We would discuss such models elsewhere.

4 Conclusion and discussion

We have studied the scenario that conformal dynamics leads to approximately R-symmetric superpotential with a metastable SUSY breaking vacuum. We have shown a simple model to realize our scenario. We can make 5D models with the same behavior. Since in our 4D scenario, metastable SUSY breaking vacua are realized by conformal dynamics, such a SUSY breaking source would be sequestered from the visible sector by conformal dynamics.

In our scenario, at a high energy scale, there would be only SUSY minimum and at low energy metastable SUSY breaking vacuum would appear. To realize the initial condition such that a metastable SUSY breaking is favored at a high energy scale, finite temperature effects would be important, because finite temperature effects might favor a metastable SUSY breaking vacuum [25].

Acknowledgement

H. A. and T. K. are supported in part by the Grand-in-Aid for Scientific Research #182496 and #17540251, respectively. T. K. is also supported in part by the Grant-in-Aid for the 21st Century COE "The Center for Diversity and Universality in Physics" from the Ministry of Education, Culture, Sports, Science and Technology of Japan.

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